

SU(3) SYMMETRY BREAKING AND OCTET BARYON POLARIZABILITIES*

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ABSTRACT

Static polarizabilities of the low-lying $\frac{1}{2}^+$ baryons are studied within the collective coordinate approach to the three flavor generalization of the Skyrme model; in particular, magnetic polarizabilities are considered. Predicted polarizabilities, which result from different treatments of the strange degrees of freedom in this model, are critically compared. Their deviations from the flavor symmetric formulations are discussed.

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At Fermilab the Σ hyperon polarizabilities will soon be measured [1, 2] and hyperon beams at CERN will provide data on the polarizabilities of other hyperons as well. In addition, a rather precise determination of the nucleon polarizabilities is available [3]. This is of great interest because the electromagnetic polarizabilities contain important information on the baryon structure[4]. Although a rather large number of theoretical work has been devoted to the nucleon electromagnetic polarizabilities (see Ref.[5] for a recent review) only quite recently the hyperon polarizabilities have been investigated. In Ref.[6] the electric and magnetic polarizabilities of the Σ_+ and Σ_- hyperons were computed within the non-relativistic quark model. A study of the hyperon polarizabilities in heavy baryon chiral perturbation theory has been reported in Ref.[7]. Within the chiral soliton models, predictions for hyperon electric polarizabilities using the $SU(3)$ collective coordinate approach have been given in Ref.[8]. Results for electric and magnetic static polarizabilities obtained within an alternative treatment of strange mesons in soliton models, the so-called bound state approach (BSA), [9], have been given recently. In this context the purpose of the present work is twofold. Firstly, we will continue the study of the polarizabilities in the $SU(3)$ collective coordinate approach to the soliton model by presenting the corresponding predictions for the static magnetic polarizabilities. Secondly, we will critically analyze and compare the results obtained within the different approaches to baryons within the $SU(3)$ Skyrme model.

Our starting point is a gauged effective chiral action

$$\Gamma = \int d^4x \left\{ \frac{f_\pi^2}{4} \text{Tr} [D_\mu U (D^\mu U)^\dagger] + \frac{1}{32\epsilon^2} \text{Tr} [U^\dagger D_\mu U, U^\dagger D_\nu U]^2 \right\} + \Gamma_{an} + \Gamma_{sb} \quad (1)$$

Here $f_\pi = 93\text{MeV}$ is the pion decay constant and ϵ is the dimensionless Skyrme parameter. Furthermore the chiral field U is the non-linear realization of the pseudoscalar octet. The covariant derivative is defined as

$$D_\mu U = \partial_\mu U + ie A_\mu [Q, U], \quad Q = \frac{1}{2} \left[\lambda_3 + \frac{1}{\sqrt{3}} \lambda_8 \right], \quad (2)$$

where A_μ is the electromagnetic field and Q the electric charge matrix. Throughout this paper we adopt Gaussian units, *i.e.* $e^2 = 1/137$. In Eq (1) Γ_{an} is the Wess-Zumino action gauged to contain the electromagnetic interaction [10] while the (gauged) symmetry breaking term Γ_{sb} [9] accounts for different masses and decay constants of the pseudoscalar fields [11]. It is convenient to order the effective action according to powers of A_μ

$$\Gamma = \Gamma^{strong} + \Gamma^{lin} + \Gamma^{quad}. \quad (3)$$

Formally we may write

$$\Gamma^{lin} = \int d^4x e A_\mu J^\mu \quad \text{and} \quad \Gamma^{quad} = - \int d^4x e^2 A_\mu G^{\mu\nu} A_\nu. \quad (4)$$

Explicit expressions for the electromagnetic current J^μ and the seagull tensor $G^{\mu\nu}$ can *e.g.* be found in ref [9]. Actually both Γ^{lin} and Γ^{quad} contribute to the baryon polarizabilities.

In second order perturbation Γ^{lin} gives rise to the so-called “dispersive” contributions while Γ^{quad} yields the so-called “seagull” contributions.

In Eq.(3) Γ^{strong} is the action in the absence of the electromagnetic field. In the soliton picture strong interaction properties of the low-lying $\frac{1}{2}^+$ and $\frac{3}{2}^+$ baryons are computed following the standard $SU(3)$ collective coordinate approach to the Skyrme model. We introduce the *ansatz*

$$U(\mathbf{r}, t) = A(t) \begin{pmatrix} c + i\boldsymbol{\tau} \cdot \hat{\mathbf{r}} s & 0 \\ 0 & 1 \end{pmatrix} A^\dagger(t) \quad (5)$$

for the chiral field. Here we have employed the abbreviations $c = \cos F(r)$ and $s = \sin F(r)$ where $F(r)$ is the chiral angle which parametrizes the soliton. The collective rotation matrix $A(t)$ is $SU(3)$ valued. Substituting the configuration (5) into Γ^{strong} yields (upon canonical quantization of A) the collective Hamiltonian. Its eigenfunctions and eigenvalues are identified as the baryon wavefunctions $\Psi_B(A) = \langle B|A \rangle$ and masses m_B . Due to the symmetry breaking terms in Γ_{sb} this Hamiltonian is obviously not $SU(3)$ symmetric. As shown by Yabu and Ando [12] it can, however, be diagonalized exactly. This diagonalization essentially amounts to admixtures of states from higher dimensional $SU(3)$ representations into the octet ($J = \frac{1}{2}$) and decouplet ($J = \frac{3}{2}$) states. This procedure, commonly known as “Rigid Rotator Approach” (RRA), has proven quite successful in describing the hyperon spectrum and static properties [13]. In ref [14] the chiral angle was allowed to adjust itself according to the flavor orientation A . This approach considers the collective rotation as slow enough to let the soliton profile react on the forces exerted by the symmetry breaking, hence the notion “Slow Rotator Approach” (SRA). In the SRA the chiral angle not only depends on the radial coordinate r but also parametrically on the flavor orientation A . In contrast to both the RRA as well as the BSA this approach has the desired feature that the meson profiles of the configuration which have their chiral field rotated maximally into the strange direction decay with the kaon mass. The comparison [14] of the predicted magnetic moments with the experimental data shows that the incorporation of symmetry breaking effects into the chiral angle is crucial to properly describe the observed deviations from U -spin symmetry¹. It is a major purpose of the present paper to compare the predictions for the magnetic polarizabilities in these approaches to the three flavor Skyrme model.

The static polarizabilities can be extracted from the shift of the particle energies in the presence of constant external electric (\mathbf{E}) and magnetic (\mathbf{B}) fields:

$$\delta M = -\frac{1}{2} \alpha \mathbf{E}^2 - \frac{1}{2} \beta \mathbf{B}^2. \quad (6)$$

The electric (α) and magnetic (β) polarizabilities characterize the dynamical response to

¹Similar results have been found by treating the influence of the symmetry breaking on the soliton extension at the quantum level [15].

the external electromagnetic fields. Here we will concentrate on the magnetic polarizability β , which is easily obtained from (3) by adopting

$$A^\mu = (0, -\frac{1}{2}\mathbf{r} \times \mathbf{B}) . \quad (7)$$

In analogy to Eq. (3) the Hamiltonian is expanded up to quadratic order in \mathbf{B}

$$H = H^{strong} + H^{lin} + H^{quad} . \quad (8)$$

The quadratic part yields the seagull contribution β_s . Using the ansatz Eq.(5) one obtains for $\frac{1}{2}^+$ baryons

$$\beta_s^B = \langle B | [\gamma_\pi^{(m)} D_{e,i}^2 + \gamma_K^{(m)} D_{e,\alpha}^2] | B \rangle . \quad (9)$$

These matrix elements are understood in the space of the collective coordinates with $D_{a,b} = \frac{1}{2}\text{Tr}(\lambda_a A \lambda_b A^\dagger)$ denoting the adjoint representation of the collective rotations. We have used the notation $i = 1, 2, 3$ and $\alpha = 4, 5, 6, 7$. Moreover, a sum over repeated indices is understood and $D_{e,a} = D_{3,a} + \frac{1}{\sqrt{3}} D_{8,a}$ refers to the electromagnetic direction. As discussed above, in the SRA the chiral angle depends on the flavor orientation A . Hence the spatial integrals

$$\gamma_\pi^{(m)} = -\frac{e^2}{9} \int d^3r \, r^2 s^2 \left[f_\pi^2 + \frac{1}{\epsilon^2} (F'^2 + \frac{s^2}{2r^2}) + \frac{2}{3} (f_K^2 - f_\pi^2) c (1 - D_{8,8}) \right] \quad (10)$$

$$\gamma_K^{(m)} = -\frac{e^2}{12} \int d^3r \, r^2 (1 - c) \left[f_K^2 + \frac{1}{4\epsilon^2} (F'^2 + \frac{s^2}{r^2}) + (f_K^2 - f_\pi^2) \frac{c-2}{3} (1 - D_{8,8}) \right] \quad (11)$$

have both explicit and implicit dependencies on A . This has to be taken care of when computing the matrix elements (9) in the SRA.

The dispersive contribution β_d arises from H^{lin} in (8). Choosing the z -axis along the \mathbf{B} field yields in second order perturbation

$$\beta_d^B = \frac{e^2}{2M_N^2} \sum_{B' \neq B} \frac{|\langle B | \mu_3 | B' \rangle|^2}{m_{B'} - m_B} . \quad (12)$$

Here B and B' refer to different baryon states and μ_3 is the magnetic moment operator. Its explicit expression for the present model can *e.g.* be found in Eqs.(13,15) of Ref.[14]. In order to compute the dispersive magnetic polarizability of a given baryon B we have to consider all possible states which are accessible from B by magnetic dipole transitions. The dominant contributions are expected from the lowest states with $|\Delta J| = |J - J'| = 1$ as these not only have the smallest mass differences but also sizable isovector contribution to the magnetic transitions [16]. For example, in the case of the nucleon the $N\Delta$ transition would then be dominant. In addition, on top of the ground state in a given spin-isospin channel the $SU(3)$ collective coordinate approach predicts states, which have their major support from higher dimensional representations of $SU(3)$. For example, states with

proton quantum numbers also exist in the $\overline{\mathbf{10}}$ and $\mathbf{27}$ representations. Such states also have non-vanishing magnetic dipole transitions to B . In Eq. (12) we have therefore included the magnetic dipole transitions to these states in both the rigid and the slow rotator approaches. Only in the limit of infinitely large symmetry breaking, when the model essentially reduces to flavor $SU(2)$, these transitions vanish.

The use of $f_K \neq f_\pi$ is essential to reproduce the experimentally observed mass differences of the low-lying $\frac{1}{2}^+$ and $\frac{3}{2}^+$ baryons [11]. For definiteness we will take $f_K = 120\text{MeV}$ and $\epsilon = 4.10$ [17] for RRA and $f_K = 118\text{MeV}$ and $\epsilon = 3.46$ [14] for SRA respectively. For the meson masses we employ $m_\pi = 138\text{MeV}$ and $m_K = 495\text{MeV}$ in both cases. In tables 1 and 2 we display the results for the dispersive contributions stemming from $|\Delta J| = 1$ and $|\Delta J| = 0$ transitions². Of course, the total dispersive magnetic polarizability is the sum of these two pieces. In these tables “1st” indicates that only that intermediate state, which has the lowest excitation energy, is included while “1st + 2nd” refers to the sum (12) being cut after the next-to-lowest state. The total contribution is obtained by including all the intermediate states with an excitation energy smaller than 3 GeV. Let us first discuss the $|\Delta J| = 1$ contributions. Here we also consider the transition $\Lambda - \Sigma_0$ although both particles have $J = \frac{1}{2}$, because these two particles are distinct by physical (isospin) quantum numbers rather than orthogonal mixtures of higher $SU(3)$ representations. Otherwise the “1st” state indeed corresponds to the observed $J = \frac{3}{2}$ baryon resonance which carries the same electrical and strangeness charges as the $J = \frac{1}{2}$ baryon under consideration. All other states (“2nd” and higher) are associated with higher $SU(3)$ excited $J = \frac{3}{2}$ states. We find that for all channels, which have a significant contribution from the “1st” state, (say, greater than one), the share carried by the excited states is almost negligible (less than 3%). Only when the “1st” contribution is small for some reason (*e.g.* it is U-spin forbidden as in the Σ_- case [18]) the “2nd” transition becomes important. In these channels the total dispersive magnetic polarizability nevertheless remains small. Basically, for all transitions the contributions from states higher than “2nd” are negligible (less than 0.5%). The RRA apparently exhibits only moderate deviations from the $SU(3)$ symmetry relations

$$\beta_d(N - \Delta) = \beta_d(\Sigma_+ - \Sigma_+^*) = \beta_d(\Xi_0 - \Xi_0^*) = \frac{4}{3}\beta_d(\Lambda - \Sigma_0^*), \quad (13)$$

which are obtained by considering only the lowest intermediate state in Eq. (12). The SRA violates these relations by as much as 50 %. Such a pattern has also been found for various other baryon properties [13]. From table 2 we observe that, as expected, the $|\Delta J| = 0$ contributions are generally quite small. Again they are only recognizable when the corresponding $|\Delta J| = 1$ transition is U-spin forbidden. Also, the contribution from the “2nd” states is important only in some particular cases (*e.g.* p , Σ_+) while all contributions from states higher than “2nd” may be discarded.

²As customary, throughout this paper all the baryon polarizabilities are expressed in units of 10^{-4} fm^3 .

The total dispersive as well as the seagull contributions are given in Table 3. There we not only compare the RRA and SRA but also quote the results from the BSA [9]. For the total dispersive part we see that the deviations from the symmetry relations (13) in the BSA and RRA are opposite with respect to those of the SRA. This result is not completely unexpected because the first two approaches inaccurately predict the magnetic moment of the Σ_+ to be slightly larger or approximately equal to that of the proton [19, 13], (in that case the symmetry relation in question would read $\mu(\Sigma_+) = \mu(p)$ [20]). As mentioned above, a major success of the SRA is the correct prediction of the pattern of the magnetic moments, especially $\mu(\Sigma_+)/\mu(p) \approx 0.85$ [14]. For the seagull contributions we again recognize that the SRA yields sizable deviations from the symmetry relations

$$\beta_s(p) = \beta_s(\Sigma_+) \quad \beta_s(n) = \beta_s(\Xi_0) \quad \beta_s(\Sigma_-) = \beta_s(\Xi_-) \quad (14)$$

while neither the RRA nor the BSA do so. In case of the SRA these deviations cause β_s to vary almost linearly with the strangeness charge, while the results from both RRA and BSA are roughly independent of strangeness. It is also somewhat surprising that while for the non-strange baryons (p, n) the predictions on β_d are comparable in the RRA and SRA they differ by a factor two in case of β_s . This indicates that strange degrees of freedom play a significant role inside the nucleon since in the infinite symmetry breaking limit, when the strange quarks are frozen out, these two approaches yield identical – $SU(2)$ – results.

Up to now we have discussed the individual contributions separately. However, the physically relevant quantity rather is the total polarizability $\beta = \beta_d + \beta_s$. The corresponding predictions for β are also given in Table 3. We observe that for the nucleon the SRA prediction is quite good because experiments favor a small positive number. The latest value quoted by the PDG [23] is $\beta(p) = 2.1 \pm 0.8 \pm 0.5$. Comparison with the prediction of the non-relativistic quark model [6] for $\beta(\Sigma_+) = 1.7$ and $\beta(\Sigma_-) = -1.7$ also favors the SRA. However, as can be seen from table 3, any treatment of the three flavor Skyrme model leads to sizable isoscalar and isotensor contributions for the magnetic polarizabilities in the Σ channel. In the Σ_0 channel the dispersive part is negative because this state dominantly couples to Λ which has a lower mass. Hence this channel is the only one where dispersive and seagull parts add coherently indicating that the Σ_0 has the largest (in magnitude) magnetic polarizability.

For completeness we also display the results for the electric seagull polarizability. The pertinent choice for the electromagnetic field is $A_\mu = (-\mathbf{E} \cdot \mathbf{r}, \mathbf{0})$. In the electric case the seagull contribution is a good approximation to the total polarizability [9, 21, 22]. In the collective treatment it is obtained from the matrix element [8]

$$\alpha_s^B = \langle B | [\gamma_\pi^{(e)} D_{e,i}^2 + \gamma_K^{(e)} D_{e,\alpha}^2] | B \rangle . \quad (15)$$

Again, these matrix elements are evaluated in the space of the collective coordinates.

Furthermore

$$\gamma_{\pi}^{(e)} = \frac{2e^2}{9} \int d^3r \, r^2 s^2 \left[f_{\pi}^2 + \frac{1}{\epsilon^2} (F'^2 + \frac{s^2}{r^2}) + \frac{2}{3} (f_K^2 - f_{\pi}^2) c (1 - D_{8,s}) \right] \quad (16)$$

$$\gamma_K^{(e)} = \frac{e^2}{6} \int d^3r \, r^2 (1 - c) \left[f_K^2 + \frac{1}{4\epsilon^2} (F'^2 + 2\frac{s^2}{r^2}) + (f_K^2 - f_{\pi}^2) \frac{c-2}{3} (1 - D_{8,s}) \right]. \quad (17)$$

Here we have omitted non-minimal photon couplings since, practically, they give no contribution to the electric polarizabilities (see footnote 3 in ref [9] for details on this issue). In table 4 the numerical results are compared to the corresponding predictions of the BSA. We observe that the SRA prediction of the electric seagull polarizability for the nucleon ($\alpha_s(p) = 11.2, \alpha_s(n) = 11.0$) agrees reasonably well with the PDG data: $\alpha(p) = 12.1 \pm 0.8 \pm 0.5$ and $\alpha(n) = 9.8_{-2.3}^{+1.9}$. Both, the RRA and the BSA yield numbers which are about twice as large. As the collective structures of the operators in (15) and (9) are identical not only the relations analogous to (14) hold in the symmetric case but also the above discussed deviations from the flavor symmetric predictions are similar for the electric and magnetic seagull contributions. To a good accuracy the seagull pieces obey $\alpha_s = -2\beta_s$ in all three treatments. This implies that the Skyrme term is only of minor importance.

We have seen that various treatments of the three flavor generalization of the Skyrme model yield quite different results for the electromagnetic polarizabilities. In particular the deviations from the $SU(3)$ symmetry relations for the dispersive parts are quite different while at least the seagull parts in the BSA and RRA are quite similar. Actually similarities between the BSA and RRA are expected from the computation of many other observables [13, 16]. Comparing especially the symmetry breaking pattern for the predicted magnetic moments of the $\frac{1}{2}^+$ baryons with the experimental data however favors the SRA. Available data on the nucleon polarizabilities tend to support this assessment. It is thus suggestive that the pattern of the electromagnetic polarizabilities of the low-lying $\frac{1}{2}^+$ baryons should follow the predictions of the SRA to the $SU(3)$ Skyrme model.

Let us finally add a word of caution concerning the quantitative results. As is well-known, neither of the three approaches discussed here correctly predicts the absolute values of the baryon magnetic moments, *e.g.* the magnetic moment of the proton is found to be (in nucleon magnetons) 1.77, 1.68 and 1.78 in the bound state, the rigid rotator and the slow rotator approaches, respectively. This is to be compared with the actual value of 2.79. This insufficiency is inherited from the $SU(2)$ Skyrme model, but it is also cured there. A recent study has shown that the moments at $\mathcal{O}(N_C)$ plus the quantum corrections at next to leading order, $\mathcal{O}(N_C^0)$, fill the gap [24]. General considerations of the $1/N_c$ -expansion show that the magnetic moment operator μ_3 acquires a multiplicative correction [25]. Since this operator crucially enters the dispersive parts of the magnetic polarizabilities (12) a change in the numerical results would not be unexpected. Similar corrections may also arise for the seagull component of the magnetic polarizabilities. The

computation of the electric polarizabilities in the two flavor models also shows that loop corrections to the corresponding $\mathcal{O}(N_C)$ seagull components are important[24]. Whether this statement carries over to $SU(3)$ remains subject to further studies.

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Table 1: Dispersive contributions β_d to the static magnetic polarizabilities corresponding to transitions between different baryons (mostly $|\Delta J| = 1$ transitions). All data are in 10^{-4}fm^3 .

	RRA			SRA		
	1st	1st + 2nd	Total	1st	1st + 2nd	Total
$N - \Delta$	4.529	4.559	4.559	5.621	5.703	5.709
$\Lambda - \Sigma_0$	4.031	4.093	4.098	3.479	3.581	3.586
$\Lambda - \Sigma_0^*$	3.835	3.875	3.877	3.237	3.318	3.322
$\Sigma_+ - \Sigma_+^*$	4.954	5.200	5.204	3.512	3.600	3.605
$\Sigma_0 - \Sigma_0^*$	0.875	1.067	1.070	0.572	0.657	0.659
$\Sigma_- - \Sigma_-^*$	0.126	0.270	0.275	0.130	0.213	0.214
$\Xi_0 - \Xi_0^*$	5.060	5.419	5.423	2.873	2.952	2.956
$\Xi_- - \Xi_-^*$	0.134	0.503	0.504	0.062	0.224	0.225

Table 2: Dispersive contributions β_d to the static magnetic polarizabilities corresponding to the $|\Delta J| = 0$ transitions. The superscript *exc* refers to $SU(3)$ excited states. All data are in 10^{-4}fm^3 .

	RRA			SRA		
	1st	1st + 2nd	Total	1st	1st + 2nd	Total
$p - p^{exc}$	0.010	0.042	0.042	0.014	0.051	0.051
$n - n^{exc}$	0.081	0.085	0.085	0.098	0.098	0.100
$\Lambda - \Lambda^{exc}$	0.017	0.017	0.017	0.051	0.051	0.051
$\Sigma_+ - \Sigma_+^{exc}$	0.023	0.079	0.080	0.005	0.030	0.030
$\Sigma_0 - \Sigma_0^{exc}$	0.086	0.110	0.111	0.021	0.033	0.033
$\Sigma_- - \Sigma_-^{exc}$	0.174	0.185	0.186	0.130	0.133	0.134
$\Xi_0 - \Xi_0^{exc}$	0.011	0.011	0.011	0.0003	0.0003	0.0003
$\Xi_- - \Xi_-^{exc}$	0.029	0.029	0.029	0.008	0.008	0.008

Table 3: Total magnetic polarizabilities as the sum of the dispersive and seagull contributions, *i.e.* $\beta = \beta_d + \beta_s$. Results are listed according to different approaches to the SU(3) Skyrme model, see text. All data are in 10^{-4}fm^3 .

	BSA [9]			RRA			SRA		
	β_d	β_s	β	β_d	β_s	β	β_d	β_s	β
p	5.6	- 8.3	- 2.7	4.6	- 10.2	- 5.6	5.8	- 5.3	0.5
n	5.6	- 8.3	- 2.7	4.6	-10.0	- 5.3	5.8	- 5.2	0.6
Λ	12.1	- 8.7	3.4	8.0	- 9.8	-1.8	7.0	- 3.2	3.6
Σ_+	10.4	- 9.1	1.3	5.3	- 10.7	-5.4	3.6	- 3.1	0.5
Σ_0	- 4.0	- 8.7	- 12.7	- 2.7	- 9.6	- 12.3	- 2.9	- 2.8	- 5.6
Σ_-	0.5	- 8.4	- 7.9	0.5	- 8.4	- 8.0	0.4	- 2.4	- 2.1
Ξ_0	14.0	- 9.6	4.4	5.4	-10.4	- 5.0	3.0	- 2.3	0.7
Ξ_-	1.5	- 8.7	- 7.2	0.5	- 7.7	- 7.2	0.2	- 1.7	- 1.4

Table 4: The electric polarizabilities as approximated by their seagull contributions (15) in various treatments of the SU(3) Skyrme model. All data are in 10^{-4}fm^3 .

	BSA [9]	RRA	SRA
p	17.3	20.9	11.2
n	17.3	20.5	11.0
Λ	18.1	20.1	7.0
Σ_+	18.1	22.0	6.6
Σ_0	18.8	19.7	5.9
Σ_-	17.4	17.3	5.1
Ξ_0	19.9	21.3	4.9
Ξ_-	18.0	15.7	3.5